# Parton content of the real photon: astrophysical implications

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We possess convincing experimental evidence for the fact that the real photon has non-trivial parton structure. On the other hand, interactions of the cosmic microwave background photons with high energy particles propagating through the Universe play an important role in astrophysics. In this context, to invoke the parton content could be convenient for calculations of the probabilities of different processes involving these photons. As an example, the cross section of inclusive resonant $W^+$ boson production in the reaction $\nu\gamma \to W^+X$ is calculated by using the parton language. Neutrino–photon deep inelastic scattering is considered.

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## I. INTRODUCTION

The parton model, according to which hadrons consist of quarks antiquarks and gluons (partons), bound together in different ways, has been very successful in reproducing experiment. This provides a relatively explicit and transparent technique for the description of high energy particle interactions. The distributions of partons inside hadrons are characterized by the structure functions satisfying the Dokshitzer–Gribov–Lipatov–Altarelli–Parisi (DGLAP) equations [1, 2, 3] or ones that are basically similar. Such a function is the probability density of finding a parton in a hadron carrying a fraction of the total hadron's momentum. Numerical solutions of the equations are in a remarkable agreement with experimental measurements, especially for the nucleon [4].

Photons being involved in high energy interactions are also able to manifest hadronic structure. One can intuitively comprehend this since the photon directly couples to quarks and therefore may split into quark–antiquark pairs. The parton contributions in two-photon processes and some crucial peculiarities of the kinematic behavior of the photon structure function have been described by Walsh and Zerwas [5]. The first work in studying quantum-chromodynamics corrections to the naive point-like structure of the photon belongs to Witten [6]. This problem was also studied in Refs. [7, 8, 9]. Introducing the evolution equations, similar to the DGLAP ones, for photons as well as the properties of the corresponding solutions were under scrutiny, for instance, in a series of papers by Glük, Reya, Grassie and Vogt [10, 11, 12]. A formulation of high energy $\gamma p$ interactions taking into account the hadronic properties of the photon was proposed in Ref. [13].

Today, we possess convincing experimental evidence for the fact that the real photon has non-trivial parton structure [14].

On the other hand, the cosmic microwave background (CMB) photons may play an important role in the formation of cosmic rays (CR). One of the brightest representatives is the Greisen–Zatsepin–Kuzmin (GZK) limit on the energy of CR [15, 16]. For example, protons of energies of over about $10^{20}$ eV would be decelerated by interaction with the CMB photons, mostly due to resonant pion production, $p\gamma \to \Delta^+ \to p\pi^0(n\pi^+)$. Other interesting processes, the $\nu\gamma$ reactions and their possible astrophysical implications, were extensively discussed in the literature (see, *e.g.*, Refs. [17, 18, 19, 20, 21, 22, 23, 24] and the references cited therein). In this context, to invoke the parton content of the real photon could be convenient for calculations of the probabilities of such processes. Here, we attempt to show the example of $W^+$ boson production in the $\nu\gamma$ scattering which may have important consequences for astrophysics [17]. Studying this reaction could also provide a test of the universality of the parton distribution functions of the photon.

## II. NEUTRINO–PHOTON REACTIONS WITHIN PARTON MODEL

Let us first consider inclusive on-shell $W^+$ boson production in the reaction $\nu_e\gamma \to W^+X$ at the resonance region using the parton language. We will view it from the center-of-mass (CMS) frame of the $\nu_e\gamma$ system. Here, for example, a substantial fraction of the CMB photons will be of energies of about $(\varepsilon E_\nu^{lab})^{1/2}$, where $E_\nu^{lab}$ is the neutrino energy in the laboratory frame defined as the frame in which the CMB is isotropic, $\varepsilon$ is the CMB photon energy (typical value $\varepsilon \sim 10^{-3}$ eV [24]). This reaction is standardly factorized into two subprocesses: the emission of a positron by the photon and annihilation of the neutrino with the positron into $W^+$ (see Fig. 1a). Then the corresponding cross section may be written as

$$\sigma(s) = \int_0^1 \hat{\sigma}(xs) f_\gamma^e(x,s)\mathrm{d}x; \tag{1}$$

here $s$ is the total CMS energy squared ($s \simeq 4\varepsilon E_\nu^{lab}$), $f_\gamma^e(x,s)$ is the probability density function to find the

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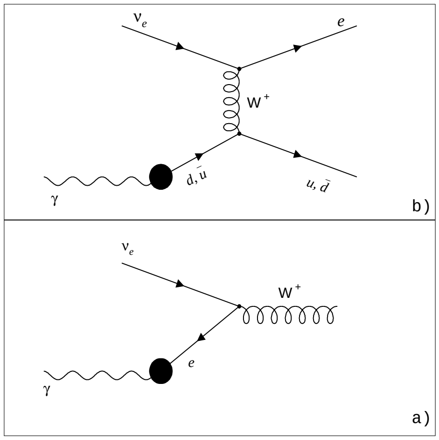


FIG. 1: Diagrams illustrating **a)** the inclusive reaction $\nu_e\gamma \to W^+X$. Neutrino annihilates with positron emitted by the photon into on-shell $W^+$; **b)** charged current neutrino scattering off quarks (antiquarks) coming from the photon. In this paper we take into account only the $u$ and $d$ quarks (antiquarks) and neglect Cabibbo–Kobayashi–Maskawa mixing.

positron in the photon carrying the fraction $x$ of the total photon's momentum, and $\hat{\sigma}(xs)$ is the cross section of the annihilation subprocess. Note that we explicitly write the $s$ dependence of the function instead of the more traditional $Q^2$ one (4-momentum transfer squared) since we deal with an $s$-channel subprocess.

In the resonance region $\hat{\sigma}(xs)$ is given by the Breit–Wigner formula [25]

$$\hat{\sigma}(xs) = 24\pi \frac{\Gamma_\mathrm{i}\Gamma}{(xs - m_W^2)^2 + m_W^2\Gamma^2}, \tag{2}$$

where $m_W$ is the mass of the $W^+$ boson, $\Gamma_\mathrm{i}$ is the partial width of the initial channel (the partial width for the decay $W^+ \to \nu_e e^+$), and $\Gamma$ is the total decay width of $W^+$. In the leading order one can find that [26]

$$\Gamma_\mathrm{i} = \frac{G_F m_W^3}{6\pi\sqrt{2}}, \quad \Gamma = 9\Gamma_\mathrm{i}, \tag{3}$$

where $G_F$ is Fermi's constant.

To determine the function $f_\gamma^e(x,s)$, we adopt the formalism given in Ref. [12]. It is fair to expect $f_\gamma^e(x,s)$ to satisfy, up to factors associated with the quark colors and fractional electric charges, the same evolution equation as the quark distributions in the photon do, provided the gluons are excluded and one takes into account only the electromagnetic interaction. Then, in the leading order we write the following equation for the positron distribution [12]:

$$\frac{\mathrm{d}\, f_\gamma^e(x,Q^2)}{\mathrm{d}\ln Q^2} = \frac{\alpha}{2\pi}k^0(x), \tag{4}$$

where $\alpha$ is the fine structure constant,
$k^0(x) = 2[x^2 + (1-x)^2]$ [27, 28]. Replacing in Eq. (4) $Q^2$ by $s$, for the reason explained above, and choosing the electron mass squared $m_e^2$ as the lower integration limit, one obtains

$$f_\gamma^e(x,s) = \frac{\alpha}{\pi}[x^2 + (1-x)^2]\ln\frac{s}{m_e^2}. \tag{5}$$

Note that the latter result is similar, for example, to the one from Ref. [29].

Substituting Eqs. (2) and (5) into Eq. (1) and performing the integration, one finally arrives at the cross section

$$\sigma(s) = \frac{8}{3}\frac{\alpha\Gamma^2}{s^3}\left[2s + \frac{s^2 - 2m_W^2(s + \Gamma^2 - m_W^2)}{\Gamma m_W} \times \left(\arctan\frac{s - m_W^2}{\Gamma m_W} + \arctan\frac{m_W}{\Gamma}\right) + (s - 2m_W^2)\ln\frac{\Gamma^2 m_W^2 + m_W^4}{\Gamma^2 m_W^2 + (s - m_W^2)^2}\right]\ln\frac{s}{m_e^2}. \tag{6}$$

The dependence of the cross section on $s$ is displayed in Fig. 2a in comparison with calculations of the closely related process $\nu_e\gamma \to W^+e^-$ carried out by Seckel [17]. Here $m_W \simeq 80.4$ GeV, $G_F \simeq 1.16\times10^{-5}$ GeV$^{-2}$, $\alpha(m_W^2) \simeq 1/128$ [30]. One can see that the values given by Eq. (6) are about two times higher than those of Ref. [17].

Let us turn now to the charged current interaction of the neutrino with the quark content of the photon (see Fig. 1b). The corresponding cross section can be obtained in the same way as it is done for neutrino–proton scattering [25]:

$$\sigma^\nu(s) = \frac{G_F^2\, s}{\pi}\left(\frac{m_W^2}{m_W^2 + Q^2}\right)^2\left(f_\gamma^q(Q^2) + \frac{1}{3}f_\gamma^{\bar{q}}(Q^2)\right), \tag{7}$$

with

$$f_\gamma^{q(\bar{q})}(Q^2) = \int_0^1 x\hat{f}_\gamma^{q(\bar{q})}(x,Q^2)\mathrm{d}x, \tag{8}$$

where $\hat{f}_\gamma^{q(\bar{q})}(x,Q^2)$ is the probability density to find a quark $q$ (antiquark $\bar{q}$) in the photon carrying the fraction $x$ of the total photon's momentum. Taking into account only the densities of the lightest quarks $u$ and $d$ from Ref. [11], we found that

$$f_\gamma^q(Q^2) = e_q^2\frac{\alpha}{2\pi}\left(\ln\frac{Q^2}{m_q^2} - \frac{3}{4}\right); \tag{9}$$

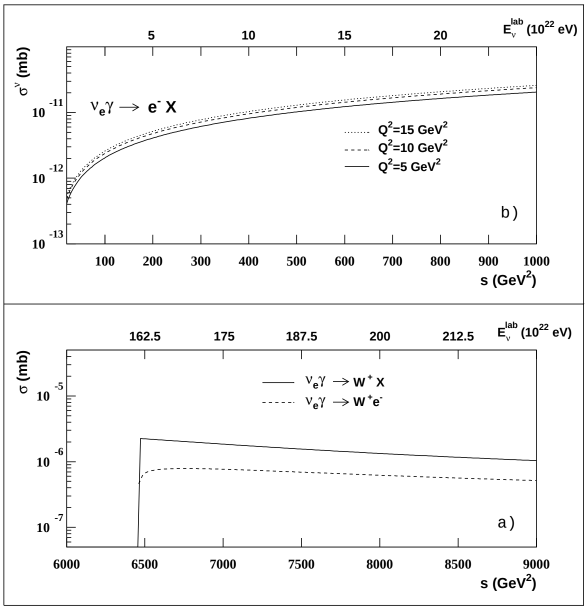


FIG. 2: **a)** Dependence of the cross section of the inclusive reaction $\nu_e\gamma \to W^+X$ on $s$ in the resonance region (solid curve). The same calculated for the closely related reaction $\nu_e\gamma \to W^+e^-$ [17] is shown by the dashed curve. **b)** Dependence of the cross section of the reaction $\nu_e\gamma \to e^-X$ on $s$ at some fixed values of $Q^2$ ($Q^2 = 5$ GeV$^2$ – solid curve, $Q^2 = 10$ GeV$^2$ – dashed curve, $Q^2 = 15$ GeV$^2$ – dotted curve). Note that the laboratory energy of the neutrino $E_\nu^{lab} \simeq s/4\varepsilon$ is calculated at $\varepsilon = 10^{-3}$ eV.

here $e_q$ and $m_q$ are the electric charge and mass of the quark $q$ respectively (for antiquarks the equation is analogous). Note that Eq. (9) is valid in the limit $m_q^2/Q^2 \ll 1$. We set $m_u = m_d = 0.2$ GeV and $\alpha = 1/137$. The dependence of the cross section thus determined on $s$ in the range 20 GeV$^2 \leq s \leq 1000$ GeV$^2$ at some values of $Q^2$ is shown in Fig. 2b.

This reaction may have interesting astrophysical implications because the struck quark may fragment into hadrons. The latter can be highly boosted and on decaying (if unstable) may produce particles with energies exceeding their GZK limit. If it occurs in the vicinity of the Earth the decay products may reach us without significant energy loss, provided the incident quark momentum pointed in the direction of the Earth. A similar idea has been proposed, for example, in Ref. [31], when photons would appear beyond the GZK limit from decays of highly boosted $\pi^0$, which, in turn, were the decay products of real $Z^0$ bosons excited in $\nu\bar{\nu}$ annihilation (the so-called "$Z$-burst" mechanism). But there are problems here, mainly associated with the origin of such high energy neutrinos, $E_\nu^{lab} \simeq m_Z^2/4\varepsilon$ (see, *e.g.*, Ref. [21]). In our case, the minimal neutrino energy required to produce hadrons is smaller than the latter one by about a factor of 400, and the corresponding cross section is also suppressed by a factor $\alpha G_F$. Anyway, one may expect that these processes were important for high energy neutrino absorption in the early Universe.

Throughout this paper we implicitly used the assumption that the parton distributions are process–independent, which has been experimentally justified for the nucleon. For example, the functions phenomenologically derived from electron–nucleon and neutrino–nucleon deep inelastic scattering data are close to each other. Using them one can correctly predict the probabilities of inclusive production of $\mu^+\mu^-$ pairs in $p\bar{p}$ collisions (Drell–Yan process) [32].

Analogously, the neutrino–photon reactions could provide an instrument for studying the universality of the parton distributions in the photon.

We have discussed only the $\nu_e\gamma$ interactions. Meanwhile, all the things we said above may be straightforwardly applied to the reactions involving the antineutrino. Likewise, heavier charged leptons can be considered. One may also include neutral current interactions in the neutrino–quark scattering. Other processes involving the CMB photons can be treated in similar way.


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